

Remark on the energy-momentum tensor in the lattice formulation of 4D $\mathcal{N} = 1$ SYM

Hiroshi Suzuki

Theoretical Research Division, RIKEN Nishina Center, Wako 2-1, Saitama 351-0198, Japan

Abstract

In a recent paper, arXiv:1209.2473 [1], we presented a possible definition of the energy-momentum tensor in the lattice formulation of the four-dimensional $\mathcal{N} = 1$ supersymmetric Yang–Mills theory (4D $\mathcal{N} = 1$ SYM), that restores the conservation law in the quantum continuum limit. In the present note, we propose a quite similar but somewhat different definition of the energy-momentum tensor (that also conserves in the continuum limit) which is superior in several aspects: In the continuum limit, the origin of the energy automatically becomes consistent with the supersymmetry and the number of renormalization constants that require a (non-perturbative) determination is reduced to two from four, the number of renormalization constants appearing in the previous construction.

Keywords: Lattice gauge theory, Supersymmetry, Energy-momentum tensor

1. Introduction

Although the lattice formulation enables one to carry out non-perturbative computations in quantum field theory, it is not straightforward to define a conserved energy-momentum tensor in the lattice formulation because the spacetime lattice explicitly breaks the translational and rotational symmetries. See, for example, Ref. [2] for a construction of the energy-momentum tensor in lattice gauge theories. In a recent paper, arXiv:1209.2473 [1], we presented a possible definition of the energy-momentum tensor in the lattice

Email address: `hsuzuki@riken.jp` (Hiroshi Suzuki)

formulation of the four-dimensional $\mathcal{N} = 1$ supersymmetric Yang–Mills theory (4D $\mathcal{N} = 1$ SYM). Our way of construction was inspired by the structure of the Ferrara–Zumino (FZ) supermultiplet [3] and the energy-momentum tensor on the lattice was defined by a renormalized, modified supersymmetry (SUSY) transformation of a renormalized SUSY current on the lattice. Then, under the assumption of parameter tuning for the supersymmetry restoration [4, 5], it was shown that the energy-momentum tensor restores the conservation law in the quantum continuum limit [1]; this can thus be a basic tool to compute physical quantities related to the energy-momentum tensor, such as the viscosity. In the present note, as a possible alternative, we propose a quite similar but somewhat different definition of the energy-momentum tensor in the lattice formulation of 4D $\mathcal{N} = 1$ SYM (that also conserves in the continuum limit). This new definition is superior in several aspects compared with the one in Ref. [1]: In the continuum limit, the origin of the energy automatically becomes consistent with SUSY and the number of renormalization constants that require a (non-perturbative) determination is reduced to two from four, the number of renormalization constants appearing in the construction in Ref. [1]. For the details of our notation, we refer the reader to Ref. [1].

2. A new definition of the energy-momentum tensor on the lattice

As Ref. [1], our starting point in the construction of the energy-momentum tensor is a renormalized SUSY Ward–Takahashi (WT) relation on the lattice:

$$\langle \partial_\mu^S \mathcal{S}_\mu(x) \mathcal{O} \rangle = \left\langle \mathcal{Z} \left[-\frac{1}{a^4} \frac{\partial}{\partial \bar{\xi}(x)} \Delta_\xi + a \mathcal{E}(x) \right] \mathcal{O} \right\rangle, \quad (2.1)$$

here and in what follows we assume that the gauge invariant operator \mathcal{O} is finite, i.e., it is already appropriately renormalized. In the left-hand of this expression, ∂_μ^S denotes the symmetric difference,

$$\partial_\mu^S f(x) \equiv \frac{1}{2a} [f(x + a\hat{\mu}) - f(x - a\hat{\mu})], \quad (2.2)$$

and $\mathcal{S}_\mu(x)$ is a renormalized Noether current associated with SUSY,

$$\mathcal{S}_\mu(x) \equiv \mathcal{Z} [\mathcal{Z}_S \mathcal{S}_\mu(x) + \mathcal{Z}_T T_\mu(x)], \quad (2.3)$$

where \mathcal{Z} , \mathcal{Z}_S and \mathcal{Z}_T are renormalization constants. In Eq. (2.3), lattice operators $S_\mu(x)$ and $T_\mu(x)$ are defined by

$$\begin{aligned} S_\mu(x) &\equiv -\sigma_{\rho\sigma}\gamma_\mu \operatorname{tr} \left\{ \psi(x) [F_{\rho\sigma}]^L(x) \right\}, \\ T_\mu(x) &\equiv 2\gamma_\nu \operatorname{tr} \left\{ \psi(x) [F_{\mu\nu}]^L(x) \right\}, \end{aligned} \quad (2.4)$$

where $[F_{\mu\nu}]^L(x)$ is a lattice transcription of the field strength; see Ref. [1] for its precise definition.

In the right-hand side of Eq. (2.1), Δ_ξ is a modified SUSY transformation on lattice variables with the localized transformation parameter $\xi(x)$,

$$\Delta_\xi \equiv \delta_\xi + \mathcal{Z}_{\text{EOM}} \delta_{F\xi}, \quad (2.5)$$

which depends on another renormalization constant \mathcal{Z}_{EOM} ; see Ref. [1] for definitions of δ_ξ and $\delta_{F\xi}$. Finally, $\mathcal{E}(x)$ in Eq. (2.1) is a dimension 11/2 operator that is given by a linear combination of renormalized operators with logarithmically divergent coefficients. In deriving Eq. (2.1) [1], we assumed that the bare gluino mass is tuned to the supersymmetric point [4, 5, 6, 7, 8] and that there is no exotic SUSY anomaly of the form of a three-fermion operator [7, 8]. The multiplicative renormalization constant \mathcal{Z} in Eq. (2.3) is chosen so that the operator $\mathcal{S}_\mu(x)$ has a finite correlation function with any renormalized operator, when the point x is far apart from the support of that operator by a finite physical distance. We note that \mathcal{Z} is at most logarithmically divergent for a dimensional reason.

In Ref. [1], a symmetric energy-momentum tensor on the lattice was defined by,

$$\mathcal{T}_{\mu\nu}(x) \equiv \frac{1}{2} [\Theta_{\mu\nu}(x) + \Theta_{\nu\mu}(x)] - c\delta_{\mu\nu} \operatorname{tr} [\bar{\psi}(x)(D + M)\psi(x)], \quad (2.6)$$

where¹

$$\Theta_{\mu\nu}(x) \equiv \frac{1}{8} (\gamma_\nu)_{\beta\alpha} \frac{\partial}{\partial \xi_\beta} [\mathcal{Z} \bar{\Delta}_\xi \mathcal{S}_\mu(x)]_\alpha, \quad (2.7)$$

and $\bar{\Delta}_\xi$ is a global modified SUSY transformation on lattice variables, that is obtained by setting $\xi(x) \rightarrow \xi$ in Eq. (2.5). c in Eq. (2.6) is a constant

¹The subscripts α and β refer to the spinor indices.

to be fixed, although it does not affect the conservation of $\mathcal{T}_{\mu\nu}(x)$. Using the SUSY WT relation (2.1), it can then be shown that the energy-momentum tensor (2.6) conserves in the continuum limit [1]. The definition through Eqs. (2.6) and (2.7) was suggested by the structure of the FZ supermultiplet [3] that the SUSY transformation of the SUSY current is basically the energy-momentum tensor.

Now, our new definition of the energy-momentum tensor on the lattice proceeds as follows: By using the renormalized SUSY current (2.3), we first define the quantity,

$$\Theta_{\mu\nu}(x; \mathcal{D}_x) \equiv -\frac{1}{8} (C^{-1}\gamma_\nu)_{\alpha\beta} a^4 \sum_{y \in \mathcal{D}_x} [\partial_\rho^S \mathcal{S}_\rho(y)]_\alpha [\mathcal{S}_\mu(x)]_\beta, \quad (2.8)$$

where \mathcal{D}_x is a hypercubic region on the lattice that contains the SUSY current $\mathcal{S}_\mu(x)$ entirely; the point x is taken as the center of the region \mathcal{D}_x so that \mathcal{D}_x is invariant under the hypercubic rotation around x . Moreover, the size of the region \mathcal{D}_x must be “macroscopic”, i.e., it must be finite in the physical unit. The definition of $\Theta_{\mu\nu}(x; \mathcal{D}_x)$ thus depends on the choice of the region \mathcal{D}_x , as its argument indicates. From this $\Theta_{\mu\nu}(x; \mathcal{D}_x)$, we define a symmetric energy-momentum tensor on the lattice, simply by symmetrizing it with respect to the indices:

$$\mathcal{T}_{\mu\nu}(x; \mathcal{D}_x) \equiv \frac{1}{2} [\Theta_{\mu\nu}(x; \mathcal{D}_x) + \Theta_{\nu\mu}(x; \mathcal{D}_x)]. \quad (2.9)$$

The idea behind the definition in Eqs. (2.8) and (2.9) is as follows: In the continuum theory, at least formally, the integral of the total divergence of the SUSY current $\int_{\mathcal{D}_x} d^4y \partial_\rho \check{S}_\rho(y)$, where the region \mathcal{D}_x contains an operator at the point x , generates the SUSY transformation,

$$-\int d^4y \frac{\delta}{\delta \bar{\xi}(y)} \delta_\xi = -\frac{\partial}{\partial \bar{\xi}} \bar{\delta}_\xi, \quad (2.10)$$

on the operator (we can read off that this is the correct normalization from the corresponding relation in the lattice theory, Eq. (2.1)). In the classical continuum theory, on the other hand, the energy-momentum tensor $\check{T}_{\mu\nu}(x)$ is given by the SUSY transformation of the SUSY current [3] as (see Ref. [1]),

$$\check{\Theta}_{\mu\nu}(x) \equiv \frac{1}{8} (C^{-1}\gamma_\nu)_{\alpha\beta} \frac{\partial}{\partial \bar{\xi}_\alpha} [\bar{\delta}_\xi \check{S}_\mu(x)]_\beta, \quad (2.11)$$

$$\check{T}_{\mu\nu}(x) = \frac{1}{2} [\check{\Theta}_{\mu\nu}(x) + \check{\Theta}_{\nu\mu}(x)] - \frac{3}{4} \delta_{\mu\nu} \text{tr} [\bar{\psi}(x) \not{D} \psi(x)]. \quad (2.12)$$

Thus one sees that the definition (2.8) is a lattice transcription of the relation expected in the continuum theory,²

$$\check{\Theta}_{\mu\nu}(x) = -\frac{1}{8}(C^{-1}\gamma_\nu)_{\alpha\beta} \int_{\mathcal{D}_x} d^4y \left[\partial_\rho \check{S}_\rho(y) \right]_\alpha \left[\check{S}_\mu(x) \right]_\beta. \quad (2.13)$$

In the classical continuum theory, the right-hand side of Eq. (2.13) is independent of the choice of the region \mathcal{D}_x because of the current conservation. In the lattice theory, however, this property is lost because the conservation law of the SUSY current is broken by $O(a)$ terms. That is, the dependence on \mathcal{D}_x in Eqs. (2.8) and (2.9) is an $O(a)$ lattice artifact and the physics in the continuum limit should not depend on the choice of the region \mathcal{D}_x .³

We note that the energy-momentum tensor (2.9) is manifestly finite, because the operator $\sum_{y \in \mathcal{D}_x} \partial_\rho^S \mathcal{S}_\rho(y)$ in Eq. (2.8), being the sum of the total divergence, does not have any overlap with the operator $\mathcal{S}_\mu(x)$; Eq. (2.9) is thus the sum of products of renormalized operators at points separated by finite physical distances.

Let us show that the lattice energy-momentum tensor $\mathcal{T}_{\mu\nu}(x; \mathcal{D}_x)$ (2.9) conserves in the continuum limit. For this, we first show the conservation of $\Theta_{\mu\nu}(x; \mathcal{D}_x)$ (2.8). From its definition (2.8) and the SUSY WT relation (2.1), we have

$$\begin{aligned} & \langle \partial_\mu^S \Theta_{\mu\nu}(x; \mathcal{D}_x) \mathcal{O} \rangle \\ &= \frac{1}{8} (C^{-1}\gamma_\nu)_{\alpha\beta} a^4 \sum_{y \in \mathcal{D}_x} \left\langle \mathcal{Z} \left[-\frac{1}{a^4} \frac{\partial}{\partial \xi(x)} \Delta_\xi + a\mathcal{E}(x) \right]_\beta \left[\partial_\rho^S \mathcal{S}_\rho(y) \right]_\alpha \mathcal{O} \right\rangle. \end{aligned} \quad (2.14)$$

Suppose now that the point x stays away from the support of the operator \mathcal{O} by a finite physical distance (we express this situation by $x \rightsquigarrow \text{supp}(\mathcal{O})$)

²On the other hand, in transcribing Eq. (2.12) to the lattice theory (2.9), we discarded the last term $-(3/4)\delta_{\mu\nu} \text{tr}[\bar{\psi}(x)\not{D}\psi(x)]$. In quantum theory, this term just acts as the zero-point energy (see Ref. [1]) and we will see below that the simple prescription (2.9) gives rise to the zero-point energy that is consistent with SUSY.

³By an argument similar to the one in what follows, it is easy to see that the difference in $\mathcal{T}_{\mu\nu}(x; \mathcal{D}_x)$ due to different choices of \mathcal{D}_x vanishes in the continuum limit, at least when the energy-momentum tensor and other renormalized operators are separated to each other by finite physical distances. This shows that, in particular, the expectation value of $\mathcal{T}_{\mu\nu}(x; \mathcal{D}_x)$ with respect to physical states becomes independent of the choice of \mathcal{D}_x in the continuum limit.

and the region \mathcal{D}_x has been chosen such that $\mathcal{D}_x \cap \text{supp}(\mathcal{O}) = \emptyset$. Then,

$$\langle \partial_\mu^S \Theta_{\mu\nu}(x; \mathcal{D}_x) \mathcal{O} \rangle = \frac{1}{8} (C^{-1} \gamma_\nu)_{\alpha\beta} \left\langle \mathcal{Z} [a\mathcal{E}(x)]_\beta \left[a^4 \sum_{y \in \mathcal{D}_x} \partial_\rho^S \mathcal{S}_\rho(y) \right]_\alpha \mathcal{O} \right\rangle, \quad \text{for } x \leftrightarrow \text{supp}(\mathcal{O}). \quad (2.15)$$

Now since the operator $\sum_{y \in \mathcal{D}_x} \partial_\rho^S \mathcal{S}_\rho(y)$ does not have any overlap with the point x , Eq. (2.15) is a correlation function of renormalized operators with no mutual overlap with an overall factor of a (in front of the operator $\mathcal{E}(x)$). Thus, Eq. (2.15) vanishes in the $a \rightarrow 0$ limit and $\Theta_{\mu\nu}(x; \mathcal{D}_x)$ conserves in the continuum limit.

Next, we consider the conservation of the anti-symmetric part of $\Theta_{\mu\nu}(x; \mathcal{D}_x)$. This follows from the same argument as Ref. [1] that uses the fact that any dimension 4 anti-symmetric rank-2 tensor in the present lattice system trivially conserves in the continuum limit.⁴ Combined above two properties of $\Theta_{\mu\nu}(x; \mathcal{D}_x)$, we have the conservation law of the symmetric part of $\Theta_{\mu\nu}(x; \mathcal{D}_x)$, Eq. (2.9). That is

$$\langle \partial_\mu^S \mathcal{T}_{\mu\nu}(x; \mathcal{D}_x) \mathcal{O} \rangle \xrightarrow{a \rightarrow 0} 0, \quad \text{for } x \leftrightarrow \text{supp}(\mathcal{O}). \quad (2.16)$$

This completes the proof of the conservation law of our lattice energy-momentum tensor (2.9).

The conservation law is a necessary condition for a Noether current to have a correct WT relation and we have observed that our energy-momentum tensor in fact restores the conservation law in the continuum limit. What is not clear as of this moment is, however, whether our energy-momentum tensor really generates correctly-normalized infinitesimal translations on renormalized operators. Technically, the presence of the dimension 11/2 (so “irrelevant”) operator in Eq. (2.1) prevents a simple argument; a further study is required on this point.⁵ Nevertheless, assuming the existence of a translational invariant quantum theory and considering the uniqueness of the conserved symmetric energy-momentum tensor in the present system [2], we believe that the answer is affirmative.

⁴To apply this argument, the operator $\Theta_{\mu\nu}(x; \mathcal{D}_x)$ must be local. This is actually the case because, under any local variation of fields, the combination $\sum_{y \in \mathcal{D}_x} \partial_\rho^S \mathcal{S}_\rho(y)$ is invariant.

⁵This remark is applied also to the construction in Ref. [1].

With the new definition, we can further show that the expectation value of the energy density vanishes in the continuum limit,

$$\langle \mathcal{T}_{00}(x; \mathcal{D}_x) \rangle = \langle \Theta_{00}(x; \mathcal{D}_x) \rangle \xrightarrow{a \rightarrow 0} 0, \quad (2.17)$$

when *periodic boundary conditions* are imposed on all the fields. This property of the energy density operator is natural from the perspective of SUSY, because Eq. (2.17) corresponds to the derivative of the supersymmetric partition function (i.e., the Witten index [13]) with respect to the temporal size of the system. In other words, Eq. (2.17) shows that the origin of the energy that is consistent with SUSY is *automatically* chosen in the continuum limit; this is a virtue of the present definition of the energy-momentum tensor compared with our previous one [1].⁶ To show Eq. (2.17), we note that $\sum_{y \in L^4} \partial_\rho^S \mathcal{S}_\rho(y) = 0$ holds under the periodic boundary conditions, where L^4 denotes the lattice of the physical size L^4 . From this,

$$\begin{aligned} \langle \Theta_{00}(x; \mathcal{D}_x) \rangle &= \frac{1}{8} (C^{-1} \gamma_0)_{\alpha\beta} a^4 \sum_{y \in L^4 - \mathcal{D}_x} \left\langle [\partial_\rho^S \mathcal{S}_\rho(y)]_\alpha [\mathcal{S}_0(x)]_\beta \right\rangle \\ &= \frac{1}{8} (C^{-1} \gamma_0)_{\alpha\beta} a^4 \sum_{y \in L^4 - \mathcal{D}_x} \left\langle \mathcal{Z} [a \mathcal{E}(y)]_\alpha [\mathcal{S}_0(x)]_\beta \right\rangle, \end{aligned} \quad (2.18)$$

where $L^4 - \mathcal{D}_x$ denotes the complement of the region \mathcal{D}_x in the lattice L^4 and we have used the SUSY WT relation (2.1) in the second equality. Then since this is a correlation function of renormalized operators with no mutual overlap with an overall factor of a , this vanishes in the continuum limit, i.e., Eq. (2.17) holds.

Our new definition in Eqs. (2.8) and (2.9) contains two unknown combinations of renormalization constants which must be determined non-perturbatively. One is the overall normalization of $\mathcal{S}_\mu(x)$, $\mathcal{Z} \mathcal{Z}_S$ and other is the ratio in $\mathcal{S}_\mu(x)$, $\mathcal{Z}_T / \mathcal{Z}_S$. See Eq. (2.3). Among these, the latter ratio $\mathcal{Z}_T / \mathcal{Z}_S$ has been non-perturbatively measured in the process to find the SUSY point in non-perturbative lattice simulations using the Wilson fermion [9, 10, 11, 12]. The

⁶We emphasize that Eq. (2.17) should hold even in a theory with the spontaneously SUSY breaking, for $a^3 \sum_{\vec{x}} \mathcal{T}_{00}(x; \mathcal{D}_x)$ to be interpreted as the energy operator appearing in the right-hand side of the SUSY algebra. For the definition of the energy density operator in a lattice formulation of the two-dimensional $\mathcal{N} = (2, 2)$ SYM [14, 15] (see also Refs. [16, 17]) that possesses the property (2.17) even before taking the $a \rightarrow 0$ limit, see Refs. [18, 19, 20, 21].

former overall normalization $\mathcal{Z}\mathcal{Z}_S$ may be determined from the expectation value of the energy operator $a^3 \sum_{\vec{x}} \mathcal{T}_{00}(x; \mathcal{D}_x)$ in a certain reference (e.g., one-particle) state. Thus, the determination of unknown constants is much simpler than our previous construction in Ref. [1] that requires the determination of other two unknown constants, \mathcal{Z}_{EOM} in Eq. (2.5) and c in Eq. (2.6). This point can be a great advantage in practical applications.

On the other hand, the new definition has an $O(a)$ ambiguity associated with the choice of the region \mathcal{D}_x in Eq. (2.8) and this ambiguity can be a possible source of the systematic error. Also, since the energy-momentum tensor is defined by the product of two SUSY currents at different points as Eq. (2.8), the application requires the computation of correlation functions with the number of arguments as twice as large compared with the correlation function of the energy-momentum tensor (e.g., one defined in Ref. [1]). Only an implementation of the present construction in actual numerical simulations will answer whether there is a real payoff or not.

We believe that the basic idea on the construction of a lattice energy-momentum tensor in the present note (and in Ref. [1]) is applicable to more general 4D supersymmetric models. For our argument on the conservation law in the continuum limit to hold, however, one has to carry out parameter fine tuning of sufficiently many numbers that ensures the SUSY WT relation (2.1). If such fine tuning is feasible for the model under consideration, our idea to construct a lattice energy-momentum tensor from the SUSY current will be useful to study physical questions in supersymmetric models, such as the spontaneous SUSY breaking, the mass and the decay constant of the pseudo Nambu–Goldstone boson associated with the (classical) dilatation invariance.

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